

Can heavy neutrinos dominate neutrinoless double beta decay?

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Abstract

Neutrinoless double beta decay can be induced by various mechanisms, involving both light and heavy exchange particles such as light Majorana neutrinos, nearly-sterile neutrinos, R-parity violating supersymmetric particles, to name a few. It has been pointed out that, typically, the contribution due to the exchange of heavy particles between the two nuclei, with masses much larger than the typical exchange momentum ~ 100 MeV, is subdominant with respect to the light neutrino one. In fact, the former is severely constrained if the same interaction induces light neutrino masses. A relevant exception to this generic conclusion is when the contribution to the light neutrino masses cancels out. Here, we focus on this case, specifically in the context of seesaw models. We perform a general analysis without restricting the study to any particular region of the parameter space, although interesting limits associated to inverse and extended seesaw-like models are discussed in more detail. It turns out that, once the relevant experimental constraints and one-loop corrections to neutrino masses are taken into account, the heavy neutrinos can dominate the process only in one of those two limits. For the inverse seesaw we find a very constrained allowed region of the parameter space, with heavy neutrino masses around 5 GeV. In general, the extended seesaw allows for a larger region, with a very hierarchical heavy spectrum which has neutrinos above and below ~ 100 MeV.

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I. INTRODUCTION

The existence of neutrino masses, strongly supported by neutrino oscillation experiments, is the first experimental evidence of physics beyond the Standard Model (SM). Furthermore, the fact that neutrino masses are smaller than the mass of the other SM fermions by several orders of magnitude calls for a “natural” New Physics (NP) explanation. Most of the models, including the very popular seesaw [1–4] ones, assume that the lepton number is not a conserved symmetry and light neutrinos are Majorana particles. An interesting experimental window to search for NP gets opened: lepton number violation processes, highly suppressed in the SM, among which neutrinoless double beta decay ($0\nu\beta\beta$ decay) experiments are the most promising. In combination with the information coming from neutrino oscillation experiments, absolute neutrino mass experiments, precision measurements, and cosmology, $0\nu\beta\beta$ decay experiments can give us precious clues in order to identify the mechanism responsible for the neutrino mass generation and provide a complementary way to look for NP, possibly not otherwise accessible at the LHC.

Although NP is necessary in order to have $0\nu\beta\beta$ decay, its effects are indirect as the light neutrinos generically dominate the process in most of the models. The key point is the fact that the light neutrino contribution and the NP one are usually related through the neutrino masses: therefore the NP contribution, suppressed by being short-range, is constrained and is generally subdominant [5]. The question if a measurable direct contribution to the $0\nu\beta\beta$ decay rate coming from NP is theoretically and phenomenologically viable is thus very interesting. This question has been addressed recently in Refs. [6, 7] in the context of different type I seesaw models. In these publications, a relevant exception to the argument above has been pointed out: the case in which the tree level light neutrino contribution, induced by NP, partially cancel. In this case, it is found that the direct NP contribution to the process is indeed relevant and can be as large as current bounds. However, no detailed discussion about the correlation among the light and heavy contributions once the one-loop corrections to neutrino masses are included in the analysis is given. The main goal of this work is to analyze to what extent this is possible in the general framework of type I seesaw models, once the relevant one-loop corrections and experimental constraints are considered.

We will first very briefly review the aspects of the $0\nu\beta\beta$ decay phenomenology relevant for our analysis. Considering a general parameterization of the neutrino mass matrix and

without restricting the analysis to any region of the parameter space, we will then study under what conditions the light neutrino contribution can be cancelled at tree level. We will include the one-loop corrections and study if the heavy neutrinos can give a dominant and measurable, i.e. at reach in the next-to-next $0\nu\beta\beta$ decay experiments, contribution to the process. We will show that, even when the tree level cancellation takes place, the light and heavy contributions are not completely decoupled once the one-loop corrections are included in the study and a dominant heavy contribution may occur only in specific regions of the parameter space.

This paper is organized as follows. In section II, we briefly review the $0\nu\beta\beta$ decay phenomenology in the general context of seesaw models, introducing the notation and the Nuclear Matrix Elements (NME) we will use. In section III, the parameterization of the neutrino mass matrix is presented, distinguishing some relevant limits and their relation with well known models, such as the inverse and extended seesaw ones. Section IV is devoted to the study of the cancellation condition of the light neutrino contribution and its tree level consequences on the heavy neutrino sector. Section V is dedicated to the analysis of the relevant corrections: higher order corrections to the seesaw expansion and one-loop corrections. The combined analysis of the $0\nu\beta\beta$ decay phenomenology, when these corrections and the relevant experimental constraints are taken into account, is presented in section V. Finally, in section VI we draw our conclusions.

II. DOMINANT HEAVY NEUTRINO CONTRIBUTION TO $0\nu\beta\beta$ DECAY?

Generically, in the context of seesaw models, the contribution to neutrinoless double beta decay from NP at scales heavier than the exchanged momentum (~ 100 MeV) can not be dominant and the light neutrinos typically dominate the process, as shown in Ref. [5].

Let us very briefly review how the above result is obtained and the possible exceptions. Following the notation in Ref. [5], and restricting the study to type I seesaw models, the $0\nu\beta\beta$ decay rate can be written as

$$\frac{\Gamma_{0\nu\beta\beta}}{\ln 2} = G_{01} \left| \sum_j U_{ej}^2 \frac{m_j}{m_e} \mathcal{M}^{0\nu\beta\beta}(m_j) \right|^2, \quad (1)$$

where G_{01} is a well known kinematic factor, U is the unitary matrix which diagonalizes the complete neutrino mass matrix both for active and sterile neutrinos, m_j are the corre-

sponding eigenvalues, i.e., the neutrino masses, and $\mathcal{M}^{0\nu\beta\beta}$ are the Nuclear Matrix Elements (NME) associated to the process. The sum should be made over all the neutrino masses, including the heavy ones.

The NME can be computed using different methods, the main two being the quasi-particle random phase approximation (QRPA) [8, 9] and the interacting shell model (ISM) [10, 11]. In this work we will make use of the NME data presented in Ref. [5] and available in Ref. [12]. They were computed for different nuclei in the context of the ISM as a function of the neutrino mass, something very convenient for our analysis. We use a notation in which the dependence on the neutrino propagator is included on $\mathcal{M}^{0\nu\beta\beta}(m_j)$, in contrast with the notation usually adopted in the literature where the propagator is expanded to factorize the mass dependence. In Fig. 1 of Ref. [5] the NME dependence on the mass of the neutrino mediating the process is depicted, showing two different regions separated by the scale of the process ~ 100 MeV:

- Below the $0\nu\beta\beta$ scale, the NME reach their maximum value and are mainly independent of the neutrino mass. For $m_i \ll 100$ MeV: $\mathcal{M}^{0\nu\beta\beta}(m_i) = \mathcal{M}^{0\nu\beta\beta}(0)$.
- The NME corresponding to neutrinos much heavier than 100 MeV are suppressed with the heavy neutrino masses and scale as $\mathcal{M}^{0\nu\beta\beta}(m_I) \propto 1/m_I^2$.

This behavior of the NME, showing two clearly different regimes, can be easily understood expanding the propagator of the neutrino mediating the process. The transition region around 100 MeV is well described in Fig. 1 of Ref. [5] since no assumptions have been made on the neutrino masses in the NME computation.

We can distinguish the following two contributions to the $0\nu\beta\beta$ decay amplitude:

$$A \propto \sum_{i=1}^3 m_i U_{ei}^2 \mathcal{M}^{0\nu\beta\beta}(m_i) + \sum_I^{\text{extra}} m_I U_{eI}^2 \mathcal{M}^{0\nu\beta\beta}(m_I), \quad (2)$$

the first term corresponding to the mostly-active neutrino contribution and the following to the extra states of the model. Here and throughout the text, we use capital letters to denote the mass index of the mostly-sterile states and lowercase letters for that of the mostly-active states.

On the other hand, since a Majorana mass coupling for the active neutrinos is forbidden by the gauge symmetry, the diagonalization of the complete mass matrix leads to the following

relation:

$$\sum_{i=1}^3 m_i U_{ei}^2 + \sum_I^{\text{extra}} m_I U_{eI}^2 = 0. \quad (3)$$

This equation, which relates the light and extra degrees of freedom of the model, should be always fulfilled at tree level and plays a fundamental role in the phenomenology of $0\nu\beta\beta$ decay.

For extra states with all the masses well above 100 MeV, using the relation given in Eq. (3), the contribution to $0\nu\beta\beta$ decay in Eq. (2) can be recast as

$$\begin{aligned} A &\propto - \sum_I^{\text{heavy}} m_I U_{eI}^2 (\mathcal{M}^{0\nu\beta\beta}(0) - \mathcal{M}^{0\nu\beta\beta}(m_I)) \\ &\approx - \sum_I^{\text{heavy}} m_I U_{eI}^2 \mathcal{M}^{0\nu\beta\beta}(0) = \sum_{i=1}^3 m_i U_{ei}^2 \mathcal{M}^{0\nu\beta\beta}(0), \end{aligned} \quad (4)$$

where we have used the fact that $\mathcal{M}^{0\nu\beta\beta}(0) \gg \mathcal{M}^{0\nu\beta\beta}(m_I)$. The contribution from the light active neutrinos thus dominates. A similar argument applies to models which implement the Type II and Type III seesaw [5], and more generically to models in which heavy sterile neutrino mixing with ν_e is introduced. As sterile neutrinos contribute to light neutrino masses, $\sum_I m_I U_{eI}^2$ is constrained by the value of the light neutrino masses while their contribution to $0\nu\beta\beta$ decay is suppressed by $\mathcal{M}^{0\nu\beta\beta}(m_I)$ making it subdominant, at least if fine tuning is not invoked as we will see in the following.

These considerations apply generically to models with extra sterile neutrinos but there are some notable exceptions.

- The case of extra states below and above 100 MeV. In this case Eq. (2) can be rewritten as

$$\begin{aligned} A &\propto \left(\sum_{i=1}^3 m_i U_{ei}^2 + \sum_I^{\text{light}} m_I U_{eI}^2 \right) \mathcal{M}^{0\nu\beta\beta}(0) + \sum_I^{\text{heavy}} m_I U_{eI}^2 \mathcal{M}^{0\nu\beta\beta}(m_I) \\ &\approx \left(\sum_{i=1}^3 m_i U_{ei}^2 + \sum_I^{\text{light}} m_I U_{eI}^2 \right) \mathcal{M}^{0\nu\beta\beta}(0), \end{aligned} \quad (5)$$

and the new states below 100 MeV may give the dominant contribution. Notice that if all the extra states are below the $0\nu\beta\beta$ scale the cancellation driven by Eq. (3) forbids the process. The same behaviour as in this type-I seesaw realization, with extra states

below and above the $0\nu\beta\beta$ scale, can take place in mixed type-I/type-II (type-I/type-III) scenario. In these scenarios NP above the $0\nu\beta\beta$ scale, such heavy sterile neutrinos or heavy triplets in the just mentioned cases, is needed to avoid the cancellation, while the extra “light” sterile neutrinos may give an important contribution to neutrino less double beta decay [5].

- Additional contributions to neutrino masses. In this case the mass relation becomes

$$\sum_{i=1}^3 m_i U_{ei}^2 + \sum_I^{\text{heavy}} m_I U_{eI}^2 = m_{LL}, \quad (6)$$

where m_{LL} is an effective Majorana mass term generated for the active neutrinos by some other mechanism. m_{LL} and $\sum_I^{\text{heavy}} m_I U_{eI}^2$ could be very large and cancel nearly exactly, keeping light neutrino masses under control. In this way, even with the contribution to $0\nu\beta\beta$ of the heavy states being weighted by the corresponding NME, a dominant effect could arise. However, it would have to overcome the suppression coming from the NME ($\mathcal{M}^{0\nu\beta\beta}(0)/\mathcal{M}^{0\nu\beta\beta}(m_I) \gg 1$) and a very high level of cancellation among $\sum_I^{\text{heavy}} m_I U_{eI}^2$ and m_{LL} in Eq. (6) would be required. This possibly implies an uncomfortably high level of fine tuning and will not be studied in this work.

- A cancellation in the light neutrino contribution: $\sum_{i=1}^3 m_i U_{ei}^2 = 0$. If this cancellation takes place, the heavy neutrinos would trivially dominate the process (at least, at tree level).

In this work we are going to focus on this last possibility. This relevant exception was studied in Refs. [6, 7, 13] and not contemplated in Ref. [5]. Of course, this cancellation in the light contribution could be obtained invoking some symmetry, and the simplest one in this context is the lepton number. The well known *inverse* [14] or *linear* [15] seesaw models, that involve small violations of the lepton number, could in principle implement this scenario. However, generating a measurable heavy neutrino contribution to the $0\nu\beta\beta$ decay is not trivial even in these models. First of all, $0\nu\beta\beta$ decay is a lepton number violating process and consequently is expected to be suppressed in this context. Moreover, the suppression of the heavy neutrino contribution with the NME ($\sim 1/m_I^2$) makes having very low scale heavy masses unavoidable in order to obtain a relevant effect. This possibility has been recently explored claiming that the heavy neutrinos could be very relevant for some

particular neutrino mass textures [6, 7]. Indeed, the main goal of this note is to check to what extent this is possible, paying special attention to the stability of the light neutrino masses under one-loop corrections and their contribution to the $0\nu\beta\beta$ decay.

III. THE MODELS

We will focus on the study of SM extensions which consist in the addition of $n+n'$ fermion gauge singlets, N_i , to the SM particle content without imposing lepton number conservation, whose Lagrangian is

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{kin}} - \frac{1}{2}\overline{N_i}M_{ij}N_j^c - (Y)_{i\alpha}\overline{N_i}\tilde{\phi}^\dagger L_\alpha + \text{h.c.}, \quad (7)$$

where \mathcal{L}_{SM} is the SM Lagrangian and \mathcal{L}_{kin} are the kinetic terms of the new fields N_i . Here, and in the rest of the paper, the subindex α denotes flavour ($\alpha = e, \mu, \tau$). Without loss of generality, the neutrino mass matrix can be expressed as

$$M_\nu = \begin{pmatrix} 0 & Y_1^T v/\sqrt{2} & \epsilon Y_2^T v/\sqrt{2} \\ Y_1 v/\sqrt{2} & \mu' & \Lambda \\ \epsilon Y_2 v/\sqrt{2} & \Lambda^T & \mu \end{pmatrix} \equiv \begin{pmatrix} 0 & m_D^T \\ m_D & M \end{pmatrix}. \quad (8)$$

Here Y_1 and Y_2 are the $n \times 3$ and $n' \times 3$ matrices that form the Dirac block m_D . The Majorana submatrix M is composed by μ' , μ and Λ : $n \times n$, $n' \times n'$ and $n \times n'$ matrices respectively. Another helpful, and widely used, basis is the one in which the Majorana submatrix for the sterile neutrinos is diagonal, which we will denote with “tilde” in the following discussion. In order to illustrate the relation between these two basis, let us consider the $n = n' = 1$ case. The Majorana submatrix M can be diagonalised as

$$OM_\nu O^T = \tilde{M}_\nu, \quad O = \begin{pmatrix} 1 & 0 \\ 0 & A \end{pmatrix}, \quad (9)$$

being A a 2×2 orthogonal matrix with rotation angle ¹

$$\tan \theta = \frac{\mu' - \mu + \sqrt{4\Lambda^2 + (\mu' - \mu)^2}}{2\Lambda}. \quad (10)$$

¹ For simplicity, all the Majorana submatrix parameters in M have been considered real.

The Majorana masses \tilde{M}_1 and \tilde{M}_2 and the Yukawa couplings \tilde{Y}_i are then given by

$$\begin{aligned}\tilde{M}_{2,1} &= \frac{1}{2} \left(\mu' + \mu \pm \sqrt{4\Lambda^2 + (\mu' - \mu)^2} \right), \\ \tilde{Y}_1 &= Y_1 \cos \theta - \epsilon Y_2 \sin \theta, \\ \tilde{Y}_2 &= Y_1 \sin \theta + \epsilon Y_2 \cos \theta.\end{aligned}\tag{11}$$

Of course, the analysis can be performed in any basis, but we will mainly work in the one in which the neutrino mass matrix is given by Eq. (8).

Notice that the mass matrix given in Eq. (8) is completely general. A particularly interesting set of models, included in Eq. (8), are those studied and summarized in Ref. [16] and which include the so-called *inverse* or *multiple* seesaw models [14, 17–19]. Lepton number is assumed to be a good global symmetry only broken in the neutrino sector through the small lepton number violation terms ϵ and/or μ and/or μ' .

The light masses are proportional to ϵ and μ as long as $\tilde{M}_i \gg \tilde{m}_D$. In this way, their smallness could be due to the one of ϵ and μ , small parameters which break the lepton symmetry. Therefore, the scale of the NP, given by Λ , can be lowered to the TeV scale or even below. This allows sizable NP effects coming from the dimension six operator which, contrary to the dimension 5 one, does not present an extra suppression with ϵ and μ as it does not violate lepton number [16]. Lepton conserving processes very suppressed in the SM as the rare decays are very promising channels to probe this kind of NP. Interesting recent analysis of $\mu \rightarrow e\gamma$, $\mu \rightarrow eee$ or $\mu \rightarrow e$ conversion in the context of low scale small lepton number violating models can be found in the literature [20–22]. In Ref. [6] these sort of models are studied in the context of the $0\nu\beta\beta$ decay using a different parameterisation based on the Casas-Ibarra one [23], which parameterises the neutrino mass matrix in terms of the light and heavy masses, the $U_{\alpha i}$ matrix and an orthogonal matrix R . Notice that the parameterisation considered here is totally general and includes the Casas-Ibarra limit in which an approximate decoupling of light and heavy sectors is assumed.

In any case, we will not restrict our study to any particular value of the parameters or, in other words, to any of the above mentioned specific limits. Nevertheless, for simplicity, we will consider the case in which only two fermion singlets ($n = n' = 1$) are added. The conclusions obtained in this work can be extended to models with larger number of right handed neutrinos, barring fine tuned cancellations. On the other hand, from neutrino oscillations we know that it is not easy to accommodate the experimental data in the region

of the parameter space between the limits: $\tilde{M}_i \gg \tilde{m}_D$ (*seesaw limit*) and $\tilde{M}_i \ll \tilde{m}_D$ (*pseudoDirac limit*). In fact, in Ref. [24] it is shown how the constraints from neutrino oscillation experiments leave those limits as the only allowed regions for $n = n' = 1$ and $\tilde{M}_1 = \tilde{M}_2$. The region of the parameter space in between is ruled out and only the pseudoDirac and seesaw limits survive. Reasonably extrapolating these results to the more general case with $\tilde{M}_1 \neq \tilde{M}_2$ studied here, leaves the seesaw limit ($\tilde{M}_i \gg \tilde{m}_D$) as the only relevant part of the parameter space in the $0\nu\beta\beta$ decay context². From now on, we will focus on the *seesaw* limit. Notice, however, that this does not necessarily mean that \tilde{M}_i have to be at the GUT or the TeV scale and can be considerably lighter [25–27].

IV. LIGHT NEUTRINO MASSES AND $0\nu\beta\beta$ DECAY

For $\tilde{M}_i \gg \tilde{m}_D$, the light neutrino mass matrix is given at tree level by

$$m_{tree} \simeq -m_D^T M^{-1} m_D \simeq \frac{v^2}{2(\Lambda^2 - \mu'\mu)} (\mu Y_1^T Y_1 + \epsilon^2 \mu' Y_2^T Y_2 - \Lambda \epsilon (Y_2^T Y_1 + Y_1^T Y_2)) , \quad (12)$$

where m_D and M are the 2×3 Dirac and 2×2 Majorana sub-matrices respectively in Eq. (8) for $n = n' = 1$. Here, we have performed the standard “seesaw” m_D/M expansion keeping the leading order terms. We will discuss later if the higher order corrections can be relevant. The contribution of the light mostly-active neutrinos to the $0\nu\beta\beta$ decay amplitude is proportional to the “ ee ” element of this effective mass matrix as

$$\begin{aligned} A_{light} &\propto \sum_{i=1}^3 m_i U_{ei}^2 \mathcal{M}^{0\nu\beta\beta}(0) \approx - (m_D^T M^{-1} m_D)_{ee} \mathcal{M}^{0\nu\beta\beta}(0) = \\ &= \frac{\mu Y_{1e}^2 + \epsilon Y_{2e} (\epsilon \mu' Y_{2e} - 2\Lambda Y_{1e})}{2(\Lambda^2 - \mu'\mu)} v^2 \mathcal{M}^{0\nu\beta\beta}(0) . \end{aligned} \quad (13)$$

Therefore, the light neutrino contribution is strictly cancelled as long as the parameters of the model satisfy the following relation

$$\mu Y_{1e}^2 + \epsilon Y_{2e} (\epsilon \mu' Y_{2e} - 2\Lambda Y_{1e}) = 0 . \quad (14)$$

This condition is fulfilled for

$$\epsilon = \mu = 0 . \quad (15)$$

² Of course, the Dirac limit will not be considered in this analysis where the $0\nu\beta\beta$ decay phenomenology is studied.

Of course, it may also be satisfied for other choices of parameters, but $\epsilon = \mu = 0$ is the most stable one under radiative corrections and higher order terms in the expansion, as we will show later. From now on we will assume that this cancellation condition is fulfilled. Obviously, setting ϵ and μ to zero leads to vanishing tree level active neutrino masses as well. However, the light neutrino masses can be generated at one-loop as we will see.

One could naively think that taking into account Eq. (3) would lead us to the same cancellation for the heavy neutrinos (see Eq. (2)), however, the dependence of the NME on m_I avoids a complete cancellation, if the heavy neutrinos are not very degenerate.

When the heavy neutrinos are above the $0\nu\beta\beta$ scale, $m_4, m_5 \gg 100$ MeV, the heavy contribution to the $0\nu\beta\beta$ decay amplitude can be approximated as

$$A_{extra} \propto \sum_I^{\text{extra}} m_I U_{eI}^2 \mathcal{M}^{0\nu\beta\beta}(m_I) \propto - (m_D^T M^{-3} m_D)_{ee} \quad (16)$$

$$= v^2 \frac{(\mu^3 + \Lambda^2(2\mu + \mu')) Y_{1e}^2 - 2\epsilon\Lambda (\Lambda^2 + \mu'^2 + \mu^2 + \mu\mu') Y_{1e}Y_{2e} + (\mu'^2 + \Lambda^2(\mu + 2\mu')) \epsilon^2 Y_{2e}^2}{2(\Lambda^2 - \mu\mu')^3},$$

which reduces to

$$A_{extra} \propto \frac{v^2 \mu' Y_{1e}^2}{2\Lambda^4}. \quad (17)$$

if the light neutrino contribution is cancelled ($\epsilon = \mu = 0$). Apparently, the above expression indicates that for large values of μ' and/or small enough Λ the heavy neutrinos may give a relevant contribution to the $0\nu\beta\beta$ decay at tree level. At this point two interesting limits of Eq. (8) arise:

- Extended seesaw limit (*ESS limit*): $\mu' \gg \Lambda$, m_D . In view of Eq. (17), this possibility appears quite appealing. This limit matches the so-called extended seesaw models [28] and corresponds to a hierarchical spectrum for the heavy neutrinos:

$$\begin{aligned} m_4 &\approx \tilde{M}_1 \approx -\Lambda^2/\mu', & U_{e4} &\approx Y_{1e}v/\sqrt{2}\Lambda, \\ m_5 &\approx \tilde{M}_2 \approx \mu', & U_{e5} &\approx Y_{1e}v/\sqrt{2}\mu', \end{aligned} \quad (18)$$

where we also show the corresponding mixing with the active neutrinos. In this regime, the lightest of the two heavy neutrinos dominates the heavy contribution. Moreover, for large enough values of μ' , m_4 becomes lighter than 100 MeV, the NME takes its maximum value and the heavy contribution to the $0\nu\beta\beta$ decay becomes independent of Λ :

$$A_{extra} \propto U_{e4}^2 m_4 \mathcal{M}^{0\nu\beta\beta}(0) \approx -\frac{Y_{1e}^2 v^2}{2\mu'} \mathcal{M}^{0\nu\beta\beta}(0). \quad (19)$$

- Inverse seesaw limit (*ISS limit*): $\Lambda \gg \mu', m_D$. This limit corresponds to one of the Minimal Flavour Violation models (MFV) studied in Ref. [16]. It is also related to the case analyzed in Ref. [6], where a different parameterization is used. In this case the heavy neutrino spectrum is quasi-degenerate, forming a quasi-Dirac pair:

$$\begin{aligned} m_4 \approx -m_5 \approx \tilde{M}_1 \approx -\tilde{M}_2 \approx \Lambda, \quad U_{e4} \approx U_{e5} \approx Y_{1e}v/2\Lambda, \\ \Delta\tilde{M} \equiv |\tilde{M}_2| - |\tilde{M}_1| \approx \mu', \end{aligned} \tag{20}$$

and we can expect lepton number violating processes such as neutrino less double beta decay to be controlled by μ' .

If all the heavy neutrinos are located below the $0\nu\beta\beta$ scale, a cancellation driven by Eq. (3) is expected at tree level, as we have already mentioned. This cancellation applies in general as long as all the heavy neutrinos are in the light regime, including the two limits distinguished above.

The approximation made in Eq. (16), $\mathcal{M}^{0\nu\beta\beta}(m_I) \propto 1/m_I^2$, does not apply if one of the heavy neutrinos (or both) is lighter than (or close to) ~ 100 MeV. However, as we have already commented, we will not restrict the analysis to any particular value of the sterile neutrino masses. This is the reason why we have made use of a numerical computation for the NME in which no approximation for the neutrino mass dependence has been considered. Notice, for instance, that the phenomenology for heavy masses around 100 MeV can be very interesting and the approximation $\mathcal{M}^{0\nu\beta\beta}(m_I) \propto 1/m_I^2$ is not very accurate in that region.

In summary, at tree level the light neutrino masses are independent of μ' (and Λ) for $\epsilon = \mu = 0$, being actually zero. However, lepton number violation processes such as $0\nu\beta\beta$ decay are sensitive to these parameters and μ' in particular. The idea behind Ref. [6, 7] is to exploit this apparent decoupling between the heavy and light contributions in order to have a measurable effect in the $0\nu\beta\beta$ decay coming from the heavy side. In the following, we will check if a heavy dominant contribution is really possible once the relevant corrections and experimental constraints are taken into account.

V. HIGHER ORDER CORRECTIONS IN THE SEESAW EXPANSION

Only the leading order in m_D/M has been considered in the expansion performed in Eq. (12). We now check if the higher order corrections may induce any relevant effects

once the tree level cancellation for the light masses takes place. The next to leading order contributions to the light neutrino masses can be written as [29]:

$$\delta m = \frac{1}{2} m_{tree} m_D^\dagger M^{-2} m_D + \frac{1}{2} \left(m_{tree} m_D^\dagger M^{-2} m_D \right)^T, \quad (21)$$

where m_{tree} is the leading order contribution given by $m_{tree} = -m_D^T M^{-1} m_D$. As they are proportional to the leading order active neutrino mass m_{tree} , they are completely irrelevant for $\mu = \epsilon = 0$. In fact, the light neutrino masses vanish for $\mu = \epsilon = 0$ at all orders in the expansion [29, 30]. Contrary to the $\mu = \epsilon = 0$ case, other choices of the parameters which satisfy the cancellation condition given in Eq. (14) are flavour dependent, giving as a result non-vanishing higher order corrections.

On the other hand, the factor $m_D^\dagger M^{-2} m_D / 2$ is nothing but the coefficient of the effective $d = 6$ operator obtained when the heavy neutrinos are integrated out of the theory [31]. This coefficient, which induces deviations from unitarity of the 3×3 lepton mixing matrix, is independent of μ' when the light neutrino cancellation ($\mu = \epsilon = 0$) takes place. Therefore, for $\mu = \epsilon = 0$, the $d = 6$ operator does not introduce any relevant μ' -dependent deviation from unitarity and μ' can escape from the corresponding constraints [32, 33], even if $\mu' \gg \Lambda$.

VI. ONE-LOOP CORRECTIONS

The one-loop corrections can be of two different types: renormalizable (i.e. the running of the parameters) or finite. In this section we will study both, starting with the renormalizable corrections. The analysis will be done after electroweak symmetry breaking (EWSB).

A. Renormalizable one-loop corrections

We are mainly interested in the running behavior of the parameters μ and ϵ , since they drive the light neutrino mass cancellation, and μ' and Λ which are the key parameters associated with the heavy contribution. We have performed the computation in the basis in which the neutrino mass matrix takes the form given by Eq. (8), in such a way that the one-loop running equations [34–36] for these parameters can be directly obtained:

$$\begin{aligned}
Q \frac{d(\epsilon Y_{2\alpha})}{dQ} &= \frac{\epsilon}{(4\pi)^2} \left[\left(T - \frac{9}{4}g^2 - \frac{3}{4}g'^2 \right) Y_{2\alpha} - \frac{3}{2} Y_{2\beta} \left((Y_l^\dagger Y_l)_{\beta\alpha} - Y_{1\beta}^* Y_{1\alpha} \right) + \frac{3}{2} \epsilon^2 Y_{2\beta} Y_{2\beta}^* Y_{2\alpha} \right], \\
Q \frac{d\mu}{dQ} &= \frac{2\epsilon}{(4\pi)^2} \left[\Lambda Y_{1\beta}^* Y_{2\beta} + \mu \epsilon Y_{2\beta}^* Y_{2\beta} \right], \\
Q \frac{d\mu'}{dQ} &= \frac{2}{(4\pi)^2} \left[\mu' Y_{1\beta}^* Y_{1\beta} + \epsilon \Lambda Y_{2\beta}^* Y_{1\beta} \right], \\
Q \frac{d\Lambda}{dQ} &= \frac{1}{(4\pi)^2} \left[\Lambda Y_{1\beta}^* Y_{1\beta} + \epsilon (\mu' Y_{1\beta}^* Y_{2\beta} + \mu Y_{2\beta}^* Y_{1\beta} + \Lambda \epsilon Y_{2\beta}^* Y_{2\beta}) \right], \tag{22}
\end{aligned}$$

with $T = \text{Tr} \left(3Y_u^\dagger Y_u + 3Y_d^\dagger Y_d + Y_l^\dagger Y_l + Y^\dagger Y \right)$ and being g and g' the $SU(2)_L$ and $U(1)_Y$ gauge coupling constants of the SM. We do not need to solve the equations to realize that the effect of the one-loop renormalizable corrections to μ and ϵ is suppressed by the tree level values of ϵ or μ . This means that the cancellation of the light active neutrino masses is stable under one-loop renormalizable corrections, as expected as a Majorana mass coupling for the active neutrinos is not allowed at tree level. For vanishing ϵ and μ at tree level, the light neutrino masses keep being zero independently of the running of the parameters (even for huge tree level inputs of μ'). This is no longer true once the finite corrections are taken into account as we show in the next subsection.

B. Finite one-loop corrections

Indeed, after EWSB, a Majorana mass for the active neutrinos is generated through finite one-loop corrections. Of course, the other Yukawa and Majorana couplings among the active and sterile neutrinos also get finite corrections, but their contribution to the light neutrino masses vanish for $\mu = \epsilon = 0$. Therefore, the dominant contribution to the light neutrino masses comes from the Majorana mass generated for the active neutrinos and is given by [37–39]

$$\delta m_{LL} = \frac{1}{(4\pi v)^2} m_D^T M \left\{ \frac{3 \ln(M^2/M_Z^2)}{M^2/M_Z^2 - 1} + \frac{\ln(M^2/M_H^2)}{M^2/M_H^2 - 1} \right\} m_D, \tag{23}$$

where m_D and M are the Dirac and Majorana sub-matrices respectively, M_Z is the mass of the Z boson and M_H the Higgs boson mass. Notice that no expansion has been performed in order to obtain this result. The structure of the correction is similar to the tree level masses but in this case no cancellation takes place for $\mu = \epsilon = 0$.

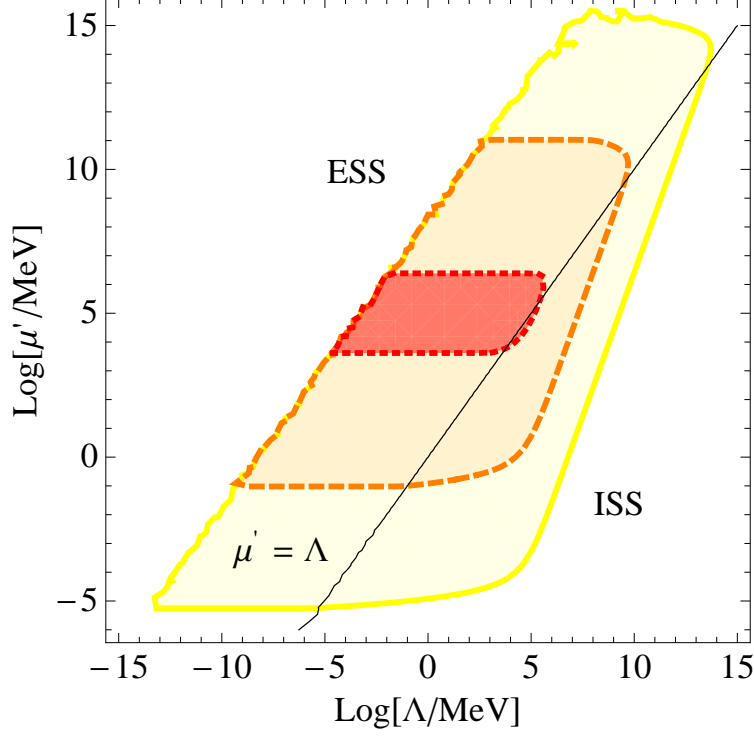


FIG. 1: The coloured region in the μ' - Λ plane corresponds to $|\delta m_{LL}(\mu = \epsilon = 0)| > 0.1$ eV. Yellow (solid), Orange (dashed) and Red (dotted) stands for $Y_{1\alpha} = 10^{-1}$, $Y_{1\alpha} = 10^{-3}$, $Y_{1\alpha} = 10^{-5}$ respectively.

Notice that the above expression is valid in any basis. In particular, Eq. (23) can be conveniently written in the $\mu = \epsilon = 0$ limit as

$$\delta m_{LL} = \frac{1}{(4\pi)^2} \frac{Y_1^T Y_1}{2} \left\{ \left(\frac{3\tilde{M}_1 \ln(\tilde{M}_1^2/M_Z^2)}{\tilde{M}_1^2/M_Z^2 - 1} + \frac{\tilde{M}_1 \ln(\tilde{M}_1^2/M_H^2)}{\tilde{M}_1^2/M_H^2 - 1} \right) \cos^2 \theta + \left(\frac{3\tilde{M}_2 \ln(\tilde{M}_2^2/M_Z^2)}{\tilde{M}_2^2/M_Z^2 - 1} + \frac{\tilde{M}_2 \ln(\tilde{M}_2^2/M_H^2)}{\tilde{M}_2^2/M_H^2 - 1} \right) \sin^2 \theta \right\}. \quad (24)$$

where the $\tilde{M}_{2,1}$ are the eigenvalues of the Majorana mass term given by Eq. (11), and θ is the rotation angle given by Eq. (10), both evaluated for $\epsilon = \mu = 0$. In Fig. 1 we show the region of the parameter space μ' - Λ given by $|\delta m_{LL}(\mu = \epsilon = 0)| > 0.1$ eV for different values of the Yukawa couplings. In order to understand better the implications of Eq. (24), we have obtained approximate expressions for two relevant limits:

- $\Lambda \gg \mu', M_H, M_Z$. We have

$$\delta m_{LL} \approx \frac{1}{(4\pi)^2} \frac{Y_1^T Y_1}{2} \frac{M_H^2 + 3M_Z^2}{\Lambda^2} \mu'. \quad (25)$$

As we have already discussed in Sec. IV, this case is included in the ISS limit and corresponds to a MFV model in which μ , ϵ and μ' are lepton number violation parameters. What we observe here is that, although the tree level light neutrino masses cancel for $\epsilon = \mu = 0$, they are generated at one loop and are proportional to the only lepton number violation parameter different from zero, μ' , as expected since the neutrino masses also violate this symmetry.

- $\mu' \gg \Lambda \gg M_H, M_Z$. In this case one finds

$$\delta m_{LL} \approx \frac{1}{(4\pi)^2} \frac{Y_1^T Y_1}{2} \left[\frac{3M_Z^2}{\mu'} \ln \left(\frac{\Lambda^4}{M_Z^4} \right) + \frac{M_H^2}{\mu'} \ln \left(\frac{\Lambda^4}{M_H^4} \right) \right]. \quad (26)$$

This case is included in the ESS limit discussed in Sec. IV. Here, the one-loop light neutrino masses depend mildly on Λ and are suppressed by μ' . Again this can be understood in terms of a lepton symmetry: μ' is suppressing the violation of Lepton number at low energies in such a way that in the limit $\mu' \rightarrow \infty$ the symmetry is completely restored in the effective theory.

In the next section we will study the phenomenological consequences of Eq. (24) in the context of the $0\nu\beta\beta$ decay without considering any expansion on the parameters. It is important to remark here that, once the tree level cancellation takes place, only one mass is generated at one-loop and at least two light masses are necessary to explain the light neutrino spectrum obtained in neutrino oscillation experiments. This is easy to solve: simply adding another fermion singlet to the model would allow to generate the necessary extra light mass. Although, for simplicity, we will keep studying the simpler case with only two extra sterile neutrinos, the conclusions extracted from the analysis can be extended to models with more than two singlets, barring fine tuned cancellations.

VII. NP DOMINANT CONTRIBUTION TO $0\nu\beta\beta$ DECAY AND ONE-LOOP NEUTRINO MASSES

Once the relevant one-loop corrections are taken into account, the Lagrangian is modified to

$$\mathcal{L} = \mathcal{L}_{\text{SM}} + \mathcal{L}_{\text{kin}} - \frac{1}{2}\overline{N}_i M_{ij} N_j^c - (\delta m_{LL})_{\alpha\beta} \overline{\nu}_{\alpha L} \nu_{\beta L}^c - (Y)_{i\alpha} \overline{N}_i \tilde{\phi}^\dagger L_\alpha + \text{h.c.} \quad (27)$$

Consequently, Eq. (3), which comes from the diagonalization of the neutrino mass matrix, is also modified to the following one-loop version:

$$\sum_i^{\text{light}} m_i U_{ei}^2 + \sum_I^{\text{extra}} m_I U_{eI}^2 = (\delta m_{LL})_{ee}. \quad (28)$$

In the case of interest, when $\epsilon = \mu = 0$, the light neutrino masses associated to the mostly-active neutrinos are determined by δm_{LL} . The tree level condition $\sum_I^{\text{extra}} m_I U_{eI}^2 = 0$ remains true at one-loop level but, as discussed in Sec. IV, heavy neutrinos could have a sizable effect to $0\nu\beta\beta$ decay thanks to the NME dependence on the heavy masses. However, at one-loop, the NP contribution to the $0\nu\beta\beta$ decay and the light neutrino masses are related as they depend on the same parameters, in particular μ' , Λ and $Y_{1\alpha}$. Their decoupling achieved at tree level does not remain true once radiative corrections are included. Consequently, the heavy parameters can not be chosen arbitrarily as to dominate $0\nu\beta\beta$ decay but are constrained by light neutrino masses, which also contribute to the process.

In principle the radiative corrections dependent on the transferred momentum p have also to be considered. These corrections are of two types: i) proportional to p or p^3 ; ii) dependent on p^2 . The first come from the W and charged Goldstone boson corrections to the neutrino propagator and vanish after performing the integration over p in the $0\nu\beta\beta$ amplitude. The second are associated to the Z and Higgs boson corrections to the propagator and are negligible in the region of heavy masses under consideration.

In the rest of the section, we will study under which conditions it may (or may not) be possible to have a dominant heavy neutrino contribution. We will pay special attention to the impact of the one-loop corrections and the experimental constraints on the parameters of the model. To illuminate the interplay among all these factors, we will first analyze the particular case of $Y_{1\alpha} = 10^{-3}$ showing our results in Fig. 2.

First of all, we are assuming that the model under consideration provides the dominant source of light neutrino masses. In principle, they should be in agreement with neutrino

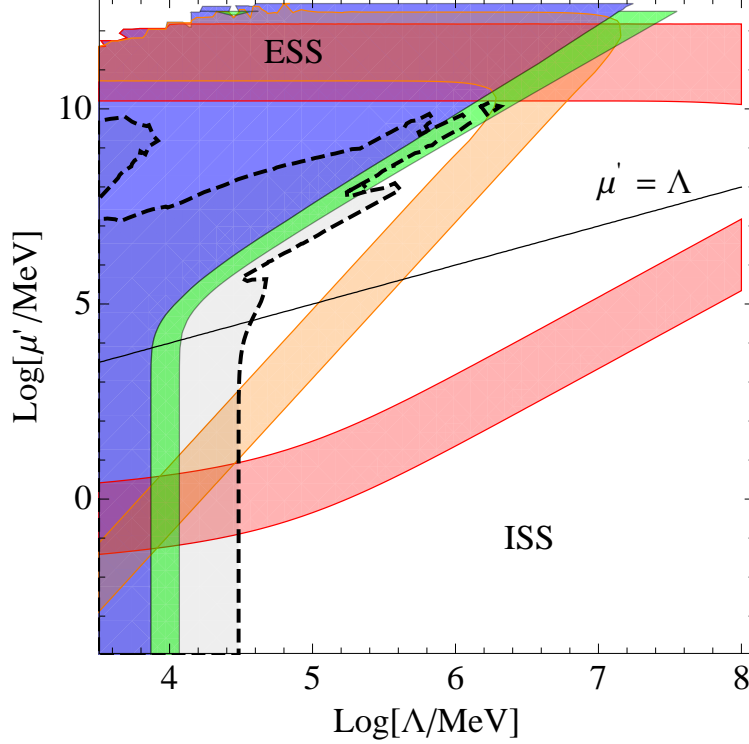


FIG. 2: Impact of one-loop corrections to $0\nu\beta\beta$ decay for $Y_{1\alpha} = 10^{-3}$. The red band is the 95%C.L. allowed region for the one-loop generated light neutrino masses bounded by cosmology and neutrino oscillations. The orange band is the 95%C.L. region of the parameter space in which the heavy neutrino contribution is between the present bound from EXO and the sensitivity of the next to next generation of $0\nu\beta\beta$ experiments. Blue (green) stands for to region in which the ratio r between the heavy and light contributions is $r > 5$ ($1 < r < 5$). The grey region inside the dashed black line is the parameter space ruled out at the 95% C.L. by the constraints on the mixing.

oscillations data but, since we are generating at one loop only one light neutrino mass we only impose a conservative lower bound on the non trivial eigenvalue of Eq. (24) given by the solar splitting, $\sqrt{\Delta m_{sol}^2}$. Moreover, the absolute neutrino mass scale experiments impose an upper bound on the same combination of parameters and as reference value we take the 95%C.L. upper bound on the light neutrino masses from cosmology [41], $m_\nu = 0.58$ eV. Since we are analyzing the case in which the tree level active neutrino masses cancel ($\epsilon = \mu = 0$), these bounds can be directly translated into bounds on μ' and Λ as a function of $Y_{1\alpha}$. They are shown in Fig. 2 as the red band. The Higgs mass, m_H , has been fixed to 125 GeV in all the calculations, as suggested by the recent LHC results [43, 44]. Notice that if no

lower bound is imposed, as it would be the case if the light neutrino masses come from some other mechanism, the outer region of the red band would not be excluded. Our conclusions remain valid also in this case, as we will discuss later. The constraint on the light neutrino masses shown in Fig. 2 can be understood analytically taking into account the approximate expressions derived in the previous section. In the ISS limit $\sqrt{\Delta m_{sol}^2} < \delta m_{LL} < 0.58$ eV scales as μ'/Λ^2 in agreement with Eq. (25). For $\mu' \gg \Lambda$, in the ESS limit, it becomes mainly independent of Λ according to Eq. (26).

The heavy neutrino contribution to $0\nu\beta\beta$ decay, given by $A_{heavy} \propto \sum_{I=4,5} U_{eI}^2 m_I \mathcal{M}^{0\nu\beta\beta}(m_I)$, can be computed diagonalizing the mass matrix in Eq. (8) and using the NME data calculated as a function of the neutrino masses [12]. The diagonalization can be easily performed in the $\epsilon = \mu = 0$ limit. This contribution has to respect the present experimental bound and, in order to be phenomenologically interesting, should be at reach of future $0\nu\beta\beta$ decay experiments such as CUORE [45], EXO [46], GERDA [47], KamLAND-Zen [48], MAJORANA [49], NEXT [50] or Super-NEMO [51]. This constraint is shown as the orange band in Fig. 2: the 95%C.L. region of the parameter space in which the heavy neutrino contribution is between the present bound from EXO [46] ($0\nu\beta\beta$ decay in ^{136}Xe), which using the corresponding shell model NME is $|m_{\beta\beta}| < 0.53$ eV, and the future (optimistic) sensitivity of the next-to-next generation of $0\nu\beta\beta$ decay experiments, taken to be $m_{\beta\beta} = 10^{-2}$ eV. The shape of the heavy contribution contour can also be easily understood from the discussion in Sec. IV. The heavy contribution scales as μ'/Λ^4 following closely Eq. (17) until, for $\mu' \gg \Lambda$, becomes independent of Λ in agreement with Eq. (19). In the ISS region, both heavy neutrinos have masses larger than the $0\nu\beta\beta$ scale and Eq. (17) holds. As expected comparing Eq. (17) and Eq. (25), in this region the slope of the heavy contribution contour is twice the $\sqrt{\Delta m_{sol}^2} < \delta m_{LL} < 0.58$ eV one. For $\mu' \gg \Lambda$ we enter the ESS limit and eventually the lightest of the two heavy masses becomes lighter than the $0\nu\beta\beta$ scale, while the heaviest one is too heavy to give a relevant contribution to the process. The “heavy” contribution is dominated thus by the sterile neutrino lighter than 100 MeV, for which the corresponding NME takes the maximum value, and is independent of Λ , see Eq. (19).

Fig. 2 also highlights the region of the parameter space for which the ratio r between the heavy and mostly-active contribution to $0\nu\beta\beta$ decay, defined as $r \equiv |A_{heavy}/A_{light}|$, is between 1 and 5 (green region) or larger than 5 (blue region). The active contribution is deter-

mined by the one-loop correction to the light neutrino masses: $A_{active} \propto (\delta m_{LL})_{ee} \mathcal{M}^{0\nu\beta\beta}(0)$. From Eqs. (17), (19) and (24), it is clear that $r \equiv |A_{heavy}/A_{light}|$ should be basically independent of the Yukawa couplings.

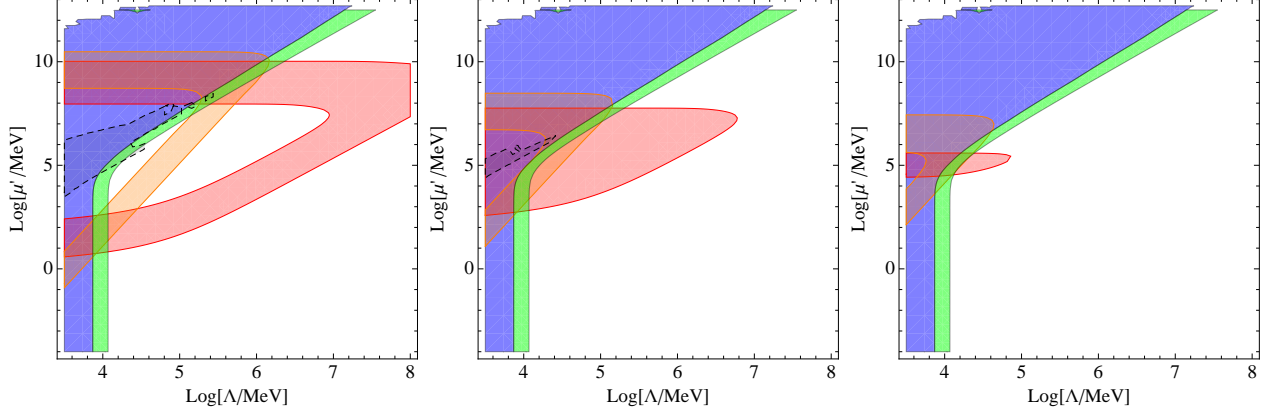


FIG. 3: Same as in Fig. 2 for $Y_{1\alpha} = 10^{-4}$ (left), $Y_{1\alpha} = 10^{-5}$ (center) and $Y_{1\alpha} = 3 \cdot 10^{-6}$ (right).

Finally, the information coming from the experiments that constrain the mixing between the active and heavy neutrinos is also included in Fig. 2. The grey region inside the dashed line is excluded at the 95% C.L. by the constraints on the mixing extracted from weak decays and summarized in [52] and non-unitarity bounds [32, 33].

As shown in Fig. 2, it is possible to have a dominant and measurable contribution from the heavy neutrinos to $0\nu\beta\beta$ decay keeping the light neutrino masses under control. Ignoring for the sake of the discussion the constraints on the heavy mixing, this takes place in the two intersections among the red, the orange and the blue regions, which lie in two interesting limits already discussed in Secs. III and IV:

- i) ISS limit: $\Lambda \gg \mu', Y_{1\alpha}v$. The heavy neutrinos are quasi-degenerate, and their contribution to the process proportional to the splitting, given by μ' . Once the constraints on the mixing, $U_{e4} \sim U_{e5} \sim Y_{1e}v/2\Lambda$, are properly taken into account the ISS limit is ruled out.
- ii) ESS limit: $\mu' \gg \Lambda, Y_{1\alpha}v$. In this case, the lightest of the extra neutrinos has mass lower than the neutrino less double beta decay exchange momentum and dominates the process.

We have chosen 10^{-3} as the input value of $Y_{1\alpha}$ in Fig. 2 as an example that allowed us to illuminate the discussion. The results for $Y_{1\alpha} = 10^{-2} - 10^{-3}$ are pretty similar, but we have

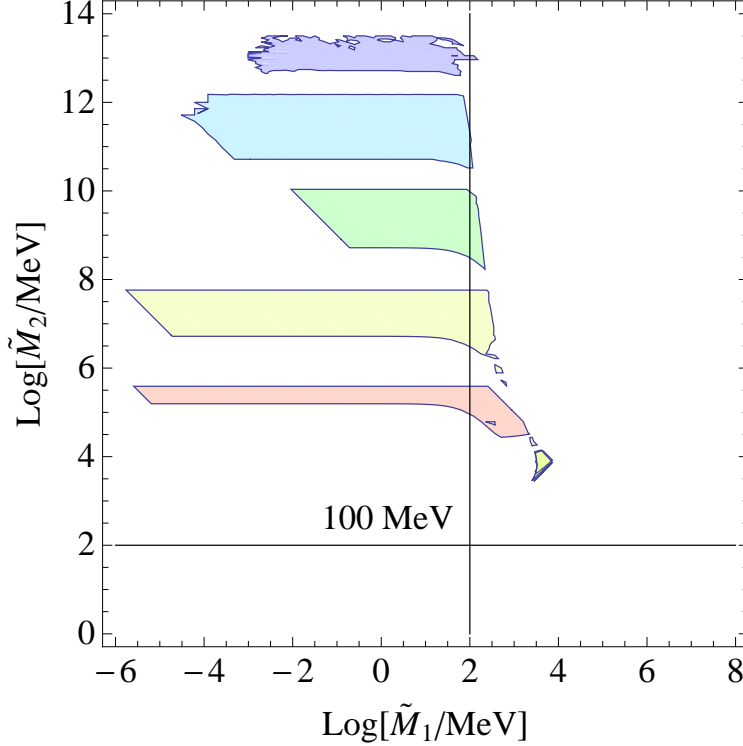


FIG. 4: Region of the parameter space, $\tilde{M}_2 \approx m_5$ vs $\tilde{M}_1 \approx m_4$, in which a dominant and measurable contribution of the heavy neutrinos, respecting bounds from neutrino oscillations, absolute neutrino mass scale experiments and weak decays, is feasible. Blue, cyan, green, yellow and red stands for $Y_{1\alpha} = 10^{-2}$, 10^{-3} , 10^{-4} , 10^{-5} and $3 \cdot 10^{-6}$ respectively. The black lines correspond to $\tilde{M}_1 = 100$ MeV and $\tilde{M}_2 = 100$ MeV.

checked that for values of the Yukawa couplings larger than 10^{-2} , a dominant contribution from the heavy neutrinos is not possible and can be at most of the same order as the one from light neutrinos. In Fig. 3 we show the plots analogous to Fig. 2 but for smaller values of the Yukawa couplings: 10^{-4} (left), 10^{-5} (center) and $3 \cdot 10^{-6}$ (right). For these values, the heavy neutrino mixing is small enough to satisfy the bounds coming from weak decays. We observe that the ratio between the light and heavy contributions is independent of the Yukawa couplings as expected. However, each of them separately depends strongly on that input. The region of the parameter space in which we have a measurable heavy contribution decreases with the Yukawa coupling, as it is also the case for the red region, in which the light neutrino masses keep being under control. We have checked that, between 10^{-6} and 10^{-8} a dominant and measurable contribution of the heavy neutrinos may still be possible,

but the light neutrino masses generated at one-loop are smaller than $\sqrt{\Delta m_{sol}^2}$. For values of the Yukawa couplings smaller than 10^{-8} the heavy contribution is too suppressed to be experimentally accessible.

The information given in Figs. 2 and 3 is summarized in Fig. 4, where we show the region of the parameter space in which a dominant and measurable contribution of the heavy neutrinos is possible respecting at the same time the bounds on heavy mixing from weak decays and non-unitarity [32, 33, 52] and keeping light neutrino masses in the region between $\sqrt{\Delta m_{solar}^2}$ and their upper bound extracted from Ref. [41]. Although the tree level cancellation for the light neutrino masses is taking place, once the one-loop corrections are taken into account, a dominant contribution from the heavy neutrinos can not occur for larger or smaller values of the Yukawa couplings than the ones shown in Fig. 4. Notice that this dominant contribution is mainly possible only in the hierarchical seesaw scenario mentioned above ($|\tilde{M}_1| \lesssim 100 \text{ MeV} \ll |\tilde{M}_2|$), where the lightest sterile neutrino gets a mass smaller than (or around) 100 MeV and dominates the process. Indeed, this result is not surprising: in Ref. [5] it was shown that, in the case in which the cancellation of the light neutrinos contribution does not occur, a hierarchical heavy spectrum like this is necessary in order to have a relevant contribution from the heavy neutrinos at tree level. We have checked in this work that this conclusion, obtained at tree level, can be extended to the case in which a cancellation of the light contribution takes place at tree level once the one-loop level corrections are included in the analysis. Nevertheless, there is an exception to these conclusions. For $Y_{1\alpha} \approx 10^{-4} - 10^{-5}$, there is still a tiny region in which the heavy contribution could dominate when the heavy neutrinos are quasi-degenerate and around 5 GeV (ISS region).

Finally, a comment is in order: the qualitative conclusions just depicted above are not affected if the lower bound on the one-loop light neutrino masses imposed here ($\delta m_{LL} > \sqrt{\Delta m_{sol}^2}$) is not assumed in the analysis. In such a case, the allowed regions in Fig. 4 become a bit larger (vertically) and a dominant heavy neutrino contribution would be still possible for $Y_{1\alpha} = 10^{-6} - 10^{-8}$. Also in this case a hierarchical spectrum with $|\tilde{M}_1| \lesssim 100 \text{ MeV} \ll |\tilde{M}_2|$ is required with the possible exception of having a quasi-degenerate spectrum with $|\tilde{M}_1| \sim |\tilde{M}_2| \sim 5 \text{ GeV}$ (in the same tiny region of the parameter space). This can be easily understood from Figs. 2-3: eliminating the lower bound on δm_{LL} would mean that the outer region of the red bands would not be forbidden any more.

VIII. SUMMARY AND CONCLUSIONS

The possibility of having a dominant contribution from “heavy” neutrinos to $0\nu\beta\beta$ decay, when a cancellation of the tree level light neutrino contribution takes place, has recently received much attention [6, 7]. In this work we have carefully analyzed this possibility in the general framework of type-I seesaw models. We have considered a general parameterization of the neutrino mass matrix which allowed us to explore the whole parameter space, identifying particularly interesting limits such as the inverse and extended seesaw models. We have shown what conditions have to be satisfied for a stable cancellation of the tree level light neutrino contribution, allowing the heavy neutrinos to dominate the process at tree level. We have studied the relevant corrections that may arise in this context. The finite one-loop corrections to the light neutrino masses turn out to be very relevant. Although logarithmic, their contribution to the $0\nu\beta\beta$ decay rate tends to dominate very easily. We have found that the heavy neutrinos can give the main contribution to the process only for a very hierarchical heavy neutrino spectrum with masses below and above the $0\nu\beta\beta$ scale ~ 100 MeV, which would match an extended seesaw like model. The heavy neutrino contribution is in fact completely dominated by the lightest sterile neutrinos with mass $\lesssim 100$ MeV, which is not suppressed by the NME. This result coincides with the general conclusions of the tree level analysis performed when no cancellation takes place [5]. Quantitatively, we have obtained that values of the Yukawa couplings between 10^{-2} and 10^{-6} (10^{-8}) are necessary, if a lower bound on the one-loop neutrino masses of $\sqrt{\delta m_{sol}^2}$ (no lower bound) is imposed in the analysis. We qualitatively agree with part of the conclusions drawn in Ref. [7]: the extended seesaw scenario might accommodate a relevant “heavy” neutrino contribution. Nevertheless, our general conclusions clarify an important detail: in Ref. [7] it was hypothesised that this may generically happen with all the heavy neutrinos above the $0\nu\beta\beta$ decay scale while we conclude that the heavy spectrum needs to contain states in both regimes, below (or close to) and above 100 MeV.

An interesting exception arises for quasi-degenerate heavy neutrinos with masses around 5 GeV which may give the dominant contribution in a tiny region of the parameter space for Yukawa couplings in the range $10^{-4} - 10^{-5}$. In agreement with Ref. [6, 7], we confirm that a relevant contribution to the $0\nu\beta\beta$ decay may come from a seesaw scenario with a quasi-degenerate heavy neutrino spectrum, which corresponds to an inverse seesaw like model.

However, we also show that this possibility is rather unlikely since it can only take place in a very particular and small region of the parameter space.

We should remark that our analysis was performed considering just two extra fermion singlets in the model. But the conclusions obtained can be safely extended to scenarios with more fermion singlets barring (more) fine tuned cancellations.

Note

During the completion of this work an analysis which considered the generation of neutrino masses at the loop-level in inverse seesaw models was presented in Ref.[53].

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